Hybrid dynamical electroweak symmetry breaking with heavy quarks and the 125 GeV Higgs

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We propose a hybrid framework for electroweak symmetry breaking (EWSB), in which the Higgs mechanism is combined with a Nambu-Jona-Lasinio mechanism. The model introduces an unconstrained (i.e., acts as "fundamental") scalar and a strongly coupled doublet of heavy quarks with a mass around 500 GeV, which forms a condensate at a compositeness scale $\Lambda \sim \mathcal{O}(1 \text{ TeV})$. This setup is matched at that scale to a tightly constrained hybrid two Higgs doublet model, where both the composite and unconstrained scalars participate in EWSB. This allows us to get a good candidate for the recently observed 125 GeV scalar which has properties very similar to the Standard Model Higgs. The heavier (mostly composite) CP-even scalar has a mass around 500 GeV, while the pseudoscalar and the charged Higgs particles have masses in the range $200 - 300$ GeV.

With the recent LHC discovery of the 125 GeV scalar particle [1], we are one step closer to understanding the mechanism of EWSB. The Standard Model (SM) Higgs mechanism lacks a fundamental explanation and has, therefore, been long questioned. Indeed, inspired from our experience with QCD, it has been speculated that the scalar responsible for EWSB may not be fundamental but, instead, some form of a fermion-antifermion bound state, which is generated dynamically by some new strong interaction at a higher scale [2].

One of the early attempts in this direction investigated the possibility of using the top-quark as the agent of Dynamical EWSB (DEWSB) via top-condensation [3], in a generalization of the Nambu-Jona-Lassinio (NJL) model [4]. However, the resulting dynamical top mass turns out to be appreciably heavier than m_t, thus making it difficult for top condensation to provide a viable picture. Moreover, top-condensate models [or NJL models where the condensing fermions have masses of $\mathcal{O}(m_t)$], require the cutoff/threshold for the new strong interactions to be many orders of magnitudes larger than m_t, i.e., of $\mathcal{O}(10^{17})$ GeV, thus resulting in a severely fine-tuned picture of DEWSB.

Nonetheless, several interesting generalizations of the top-condensate model of [3], which potentially avoid these obstacles, have been suggested. For example, one can relax the requirement that only the top-condensate is responsible for the full EWSB [5, 6], or assume that condensations of new heavier quarks and/or leptons drive EWSB [7–11]. In such scenarios the resulting low-energy (i.e., EW-scale) effective theory may contain more than a single composite Higgs doublet [6, 8–13]. Indeed, low-energy multi-Higgs models, with new heavy fermions with masses of $\mathcal{O}(500 \text{ GeV})$, are natural outcomes of a TeV-scale DEWSB scenario, since the heavy fermions are expected to be strongly coupled at the near by TeV-scale and to lead to the formation of several condensates - possibly with sub-TeV masses. On the other hand, the approach with heavy fermions has one major caveat: in "conventional" NJL models for DEWSB the typical mass of the heavy fermionic condensate, $H \sim < \bar{\Psi}\Psi >$, tends to lie in the range $m_\Psi < m_H < 2m_\Psi$ (see e.g., [3]). Thus, when $m_\Psi \sim \mathcal{O}\left(500 \; GeV\right)$, such a composite tends to be too heavy to account for the recently discovered 125 GeV Higgs-like particle.[1]

In this paper we propose an alternative solution for the TeV-scale DEWSB scenario, in which a light SM-like Higgs with a mass of $\mathcal{O}(m_W)$ emerges. The idea is to add an unconstrained scalar field at the compositeness scale (which behaves as a "fundamental" field - possibly resulting from an underlying strong dynamics, see e.g., [5]), where additional super-critical attractive 4-Fermi operators form a composite scalar sector. This strongly coupled composite-plus-fundamental sectors are matched at the compositeness scale to a hybrid "4th generation"[2] 2HDM (h4G2HDM), with one fundamental-like field (Φ_ℓ) which couples to the SM's light fermions and one auxiliary (composite) field (Φ_h) which couples to the heavy quarks.[3] The fundamental or unconstrained Higgs field is thus responsible for the mass generation of the lighter SM fermions and for the observed CKM flavor

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[1] One possible way out, which we will not consider here, is that the lightest scalar state is the pseudoscalar associated with DEWSB (see e.g., [13]), since its mass does not receive large corrections from loops of the heavy fermions and, thus, can in principle be held small without fine-tuning.

[2] The name "4th generation" is used here for convenience only and should not be confused with the minimal SM4 framework, as we explicitly involve a 2HDM. Besides, the DEWSB mechanism proposed here can be generalized to models with non-sequential heavy quarks.

[3] We note another interesting hybrid multi-Higgs model suggested in [11], in which the condensates which drive DEWSB are formed by exchanges of the fundamental Higgs.

pattern.

This h4G2HDM setup is motivated by the concept that new heavy fermions are expected to have purely dynamical masses, while a different mechanism is expected to underly the generation of mass for the lighter SM fermions. It further allows us to get a light 125 GeV SM-like Higgs, since its mass is mostly proportional to the quartic coupling of the fundamental field, which is unconstrained at the compositeness scale. That is, with this assignment, in our h4G2HDM the mass of the lightest CP-even Higgs state does not receive the usual large quantum corrections from loops of the heavy dynamical fermions, which instead feed into the quartic coupling of the composite scalar.

It should be noted that the simplest low-energy effective setup which may result from a TeV-scale DEWSB scenario is the so called SM4, i.e., the SM with a 4th sequential generation of heavy fermions and one Higgs doublet. This minimal framework has, however, several drawbacks [15] and, more importantly, it fails to account for the existence of the recently discovered 125 GeV Higgs-like particle [14]. On the other hand, a 4th generation framework with two (or more) Higgs doublets, such as the h4G2HDM discussed here, where the new heavy quarks have masses of $\mathcal{O}\,(500\text{ GeV})$, is not only consistent with the current Higgs data[4], but it can also have specific flavor structures, which may give rise to new signatures of 4th generation quarks at the LHC [15, 16], and substantially relax the current bounds on their masses [17]. Moreover, the lightest Higgs state in these class of models may be a good candidate for the recently discovered 125 GeV Higgs-like particle [18].

We assume that the 4th generation quarks are charged under some new strong interaction that dynamically breaks EW symmetry (see e.g., the "top-color" models of [5, 20]). The theory at the compositeness scale, Λ, can then be parameterized by adding to the light SM degrees of freedom the following set of strongly coupled 4-Fermi terms:

$$\begin{aligned}\mathcal{L} = \mathcal{L}_{SM}(\Lambda) \; &+ \; G_{t'}\bar{Q}'_L t'_R \bar{t}'_R Q'_L + G_{b'}\bar{Q}'_L b'_R \bar{b}'_R Q'_L \\ &+ \; G_{t'b'}\left(\bar{Q}'_L b'_R \bar{t}'^c_R i\tau_2 Q'^c_L + h.c.\right) \;, \quad (1)\end{aligned}$$

where $Q'_L = (t'_L \; b'_L)^T$ and $\mathcal{L}_{SM}(\Lambda)$ stands for the bare SM Lagrangian with a single fundamental Higgs field, Φ_ℓ, which essentially parameterizes our ignorance as for the origin of mass for the lighter fermions. It is given by:

$$\begin{aligned}\mathcal{L}_{SM}(\Lambda) &= |D_\mu \Phi_\ell|^2 + \mathcal{L}^Y_{SM}(\Lambda) - V_{SM}(\Lambda) \;, \\ \mathcal{L}^Y_{SM}(\Lambda) &= g_u^{0,ij}\bar{Q}^i_L\tilde{\Phi}_\ell u^j_R + g_d^{0,ij}\bar{Q}^i_L\Phi_\ell d^j_R + h.c. \;, \\ V_{SM}(\Lambda) &= (\mu^0_\ell)^2\Phi^\dagger_\ell\Phi_\ell + \frac{1}{2}\lambda^0_\ell\left(\Phi^\dagger_\ell\Phi_\ell\right)^2 \;, \qquad (2)\end{aligned}$$

[4] For discussions on the phenomenology of multi-Higgs 4th generation models, see e.g., [12, 15–19].

where $\tilde{\Phi} \equiv i\tau_2\Phi^\star$, $i,j = 1-3$ and we use the superscript 0 to denote bare couplings at the scale Λ .

We can reproduce the theory defined by Eq. 1 by introducing at Λ an auxiliary Higgs doublet, Φ_h, which couples only to the 4th generation quarks as follows:

$$\begin{aligned}\mathcal{L}_{q'}(\Lambda) &= g^0_{b'}\left(\bar{Q}'_L\Phi_h b'_R + h.c.\right) + g^0_{t'}\left(\bar{Q}'_L\tilde{\Phi}_h t'_R + h.c.\right) \\ &- (\mu^0_h)^2\Phi^\dagger_h\Phi_h \;. \qquad (3)\end{aligned}$$

For simplicity we did not write above the Yukawa terms for the light and for the 4th generation leptons. In particular, our results are not sensitive to the choice by which the 4th generation leptons couple to the Higgs sector; they can either couple to the fundamental Higgs or to the auxiliary field. In either case, we assume that their couplings are sub-critical and, therefore, do not play any role in DEWSB (see also discussion below).

The scalar sector of the full theory at the compositeness scale is, therefore, described by:

$$\mathcal{L}(\Lambda) = \mathcal{L}_{SM}(\Lambda) + \mathcal{L}_{q'}(\Lambda) + (\mu^0_{h\ell})^2\left(\Phi^\dagger_h\Phi_\ell + h.c.\right) \;, \quad (4)$$

where we have added a $\Phi_h - \Phi_\ell$ mixing term ($\propto \mu^0_{h\ell})^2$, which may arise e.g., from QCD-like instanton effects associated with the underlying strong dynamics (see e.g., [5, 21]) or from sub-critical couplings of the fundamental Higgs to the 4th generation quarks (see below). This term explicitly breaks the $U(1)$ Peccei-Quinn (PQ) symmetry [22], which is otherwise possessed by the model, thus avoiding the presence of a massless pseudoscalar in the spectrum. Note that, in any realistic scenario we expect $\mu_{h\ell}(\mu \sim m_W) \sim \mathcal{O}(m_W)$ and, since this is the only term which breaks the PQ symmetry, it evolves only logarithmically under the RGE so that, at the compositeness scale, we also have $\mu^0_{h\ell} \equiv \mu_{h\ell}(\mu \sim \Lambda) \sim \mathcal{O}(m_W)$. Therefore, since $\mu^0_{h/\ell} \equiv \mu_{h/\ell}(\mu \sim \Lambda) \sim \mathcal{O}(\Lambda)$, we expect $(\mu^0_{h\ell})^2/(\mu^0_h)^2 \sim \mathcal{O}(m^2_W/\Lambda^2) \ll 1$.

When the auxiliary (composite) field is integrated out at Λ, we recover the Lagrangian defined by Eq. 1, with

$$G_{t'} = \frac{(g^0_{t'})^2}{(\mu^0_h)^2} \;,\; G_{b'} = \frac{(g^0_{b'})^2}{(\mu^0_h)^2} \;,\; G_{t'b'} = -\frac{g^0_{t'}g^0_{b'}}{(\mu^0_h)^2} \;, \qquad (5)$$

plus additional interaction terms between the light Higgs and the new heavy quarks with Yukawa couplings of $\mathcal{O}\left(\frac{(\mu^0_{h\ell})^2}{(\mu^0_h)^2}\cdot g^0_{t'/b'}\right)$; since at Λ we have $(\mu^0_{h\ell})^2/(\mu^0_h)^2 \sim \mathcal{O}(m^2_W/\Lambda^2) \ll 1$, such residual $\Phi_\ell \bar{q}' q'$ terms are expected to be small and will, therefore, not participate in EWSB.

Thus, from the point of view of DEWSB, the hybrid 2HDM defined at Λ, i.e., with one fundamental and one auxiliary/composite scalar fields, is exactly equivalent to the theory defined in Eq. 1 with the strong 4-Fermi interactions of the 4th generation quarks. The auxiliary field, Φ_h, is, therefore, viewed as a composite of the form $\Phi_h \sim g_{t'} < \bar{Q}'_L t'_R > + g_{b'} < \bar{Q}'_L b'_R >$, which is responsible for EWSB and for the dynamical mass generation of the heavy quarks. At low energies the field Φ_h acquires a kinetic term as well as self interactions and the

theory behaves as a 2HDM with a structure similar to the 4G2HDM proposed in [15], i.e., one Higgs field (Φ_h) couples only to the heavy 4th generation quarks and the 2nd Higgs field (Φ_ℓ) couples to the SM quarks of the 1st-3rd generations. The mass terms μ_h, μ_ℓ receive quantum corrections, resulting in EW-scale VEV's for Φ_h and for Φ_ℓ which break the EW symmetry.

The low-energy parameters (masses and couplings) of this h4G2HDM are determined by the RGE of this model with the compositeness boundary conditions at Λ. The resulting low-energy h4G2HDM scalar potential can be written as:

$$\begin{aligned} V_{h4G2HDM}(\Phi_h,\Phi_\ell) &= \mu_\ell^2 \Phi_\ell^\dagger \Phi_\ell + \mu_h^2 \Phi_h^\dagger \Phi_h - \mu_{h\ell}^2 \left(\Phi_h^\dagger \Phi_\ell + h.c.\right) \\ &+ \frac{1}{2}\lambda_\ell \left(\Phi_\ell^\dagger \Phi_\ell\right)^2 + \frac{1}{2}\lambda_h \left(\Phi_h^\dagger \Phi_h\right)^2 + \lambda_3 \left(\Phi_h^\dagger \Phi_h\right)\left(\Phi_\ell^\dagger \Phi_\ell\right) + \lambda_4 \left(\Phi_h^\dagger \Phi_\ell\right)\left(\Phi_\ell^\dagger \Phi_h\right) \ . \end{aligned} \tag{6}$$

where all the above mass terms and quartic couplings run as a function of the energy scale μ, as dictated by the RGE for this model.[5] The stability condition for the above potential reads $\lambda_\ell, \lambda_h > 0$ and $\sqrt{\lambda_\ell \lambda_h} > -\lambda_3 - \lambda_4$. Also, we apply the compositeness boundary conditions at the scale Λ to the strong 4th generation Yukawa couplings that generate the 4-Fermi terms and to the quartic couplings involving the auxiliary field Φ_h:

$$\begin{aligned} & g_{q'}(\Lambda) \to \infty \ \ , \ \ \lambda_{h,3,4}(\Lambda) \to \infty \ , \\ & \frac{\lambda_h(\Lambda)}{g_{q'}^4(\Lambda)} \to 0 \ \ , \ \ \frac{\lambda_{3,4}(\Lambda)}{g_{q'}^2(\Lambda)} \to 0 \ , \end{aligned} \tag{7}$$

where $q' = t', b'$. Eq. 7 reflects the fact that the h4G2HDM is matched at Λ to the theory defined by Eq. 1, with the strongly coupled 4-Fermi factors derived in Eq. 5. On the other hand, the quartic coupling of the fundamental Higgs is unconstrained at Λ, so that we have $\lambda_\ell(\mu \to \Lambda) \to \lambda_\ell^{(0)}$, where $\lambda_\ell^{(0)}$ is a free parameter of the model.

The boundary conditions for μ_h, μ_ℓ and $\mu_{h\ell}$ are not required for our analysis, since we are only interested in their values at the EW-scale, which are fixed by the minimization conditions of the scalar-potential. In particular, at the minimum of the potential (i.e., at the EW-scale) we can express μ_h and μ_ℓ in terms of $\mu_{h\ell}$, $\tan\beta \equiv t_\beta = v_h/v_\ell$, $v = \sqrt{v_h^2 + v_\ell^2}$ and the quartic couplings λ_h, λ_ℓ (neglecting λ_3 and λ_4, see below):

$$\mu_\ell^2 \simeq \frac{\mu_{h\ell}^2}{t_\beta} - \frac{v^2}{2} c_\beta^2 \lambda_\ell \ \ , \ \ \mu_h^2 \simeq t_\beta \mu_{h\ell}^2 - \frac{v^2}{2} s_\beta^2 \lambda_h \ , \tag{8}$$

where $s_\beta, c_\beta = \sin\beta, \cos\beta$ and it is understood that the quartic couplings are evaluated at $\mu \sim v$, i.e., $\lambda_h = \lambda_h(\mu \sim v)$ and $\lambda_\ell = \lambda_\ell(\mu \sim v)$.

Solving the RGE of the h4G2HDM with the compositeness boundary conditions in Eq. 7, we find that $d\lambda_{3,4}/d\mu$ is proportional only to the (small) gauge couplings, so that $\lambda_3(v), \lambda_4(v) << \lambda_\ell(v), \lambda_h(v)$ and can therefore be neglected. Also, in our h4G2HDM, the RGE for λ_ℓ and for the top-quark Yukawa coupling g_t are similar to the SM RGE for these couplings. As for the RGE for λ_h and for the Yukawa couplings of the 4th generation quarks $g_{q'}(\mu)$, we can obtain a viable approximate analytic solution by neglecting the contributions from the running of the gauge couplings and the Yukawa couplings of all light fermions, as well as the Yukawa couplings of the 4th generation leptons. In particular, taking for simplicity $g_{t'} = g_{b'} \equiv g_{q'}$, the dominant (approximate) RGE in our h4G2HDM are given by:

$$\mathcal{D} g_{q'} \ \approx \ 6 g_{q'}^3 \ , \tag{9}$$

$$\mathcal{D} \lambda_h \ \approx \ 4\lambda_h \left(3\lambda_h + 6 g_{q'}^2\right) - 24 g_{q'}^4 \ , \tag{10}$$

where $\mathcal{D} \equiv 16\pi^2 \mu \frac{d}{d\mu}$. With the compositeness boundary conditions of Eq. 7, the above RGE's have a simple analytic solution:

$$g_{q'}(\mu) = \sqrt{\frac{4\pi^2}{3 ln\frac{\Lambda}{\mu}}} \ \ , \ \ \lambda_h(\mu) = \frac{4\pi^2}{3 ln\frac{\Lambda}{\mu}} \ . \tag{11}$$

Thus, using $m_{q'} = v_h g_{q'}(\mu = m_{q'})/\sqrt{2}$, we can obtain the cutoff Λ as a function of $m_{q'}$ and t_β:

$$\Lambda \approx m_{q'} \cdot exp\left(\frac{2\pi^2 \left(s_\beta v\right)^2}{3 m_{q'}^2}\right) \ . \tag{12}$$

In particular, for $m_{q'} = 400 - 600$ GeV and $t_\beta \sim \mathcal{O}(1)$, for which our low-energy h4G2HDM successfully accounts for the recently discovered 125 GeV Higgs-like particle (see below), we find that $\Lambda \sim 1 - 1.5$ TeV. Remarkably this is some fourteen orders of magnitudes smaller than the cutoff which emerges from a top-condensate scenario: $\Lambda \sim m_t \cdot exp\left(\frac{16\pi^2 v^2}{9 m_t^2}\right) \sim 10^{17}$ GeV, i.e., obtained by solving the SM-like RGE for g_t: $\mathcal{D} g_t \approx \frac{9}{2} g_t^3$. Thus, introduction of a heavy quark doublet

[5] The most general 2HDM potential also includes the quartic couplings $\lambda_{5,6,7}$ [23], which, in our h4G2HDM, are absent at any scale.

significantly alleviates the inherent fine-tuning (i.e., hierarchy) problem that afflicts the DEWSB models where the condensing fermions have masses $\lesssim 200$ GeV.[6]

The physical scalar masses are given by:

$$m_A^2 = m_{H^+}^2 = \frac{\mu_{h\ell}^2}{s_\beta c_\beta} , \tag{13}$$

$$m_{h,H}^2 = \frac{1}{2}\left(m_1^2 + m_2^2 \mp \sqrt{(m_1^2 - m_2^2)^2 + 4\mu_{h\ell}^4}\right) , \tag{14}$$

where (see also Eq. 8)

$$m_1^2 \simeq \mu_h^2 + \frac{3}{2} s_\beta^2 v^2 \lambda_h \simeq t_\beta \mu_{h\ell}^2 + s_\beta^2 v^2 \lambda_h , \tag{15}$$

$$m_2^2 \simeq \mu_\ell^2 + \frac{3}{2} c_\beta^2 v^2 \lambda_\ell \simeq \frac{\mu_{h\ell}^2}{t_\beta} + c_\beta^2 v^2 \lambda_\ell , \tag{16}$$

and a Higgs mixing angle:

$$\tan 2\alpha \simeq \left(\cot 2\beta - \frac{v^2}{2\mu_{h\ell}^2} \cdot \left(s_\beta^2 \lambda_h - c_\beta^2 \lambda_\ell\right)\right)^{-1} , \tag{17}$$

defined by:

$$\begin{aligned} h &= \cos\alpha \cdot \mathrm{Re}(\Phi_\ell^0) - \sin\alpha \cdot \mathrm{Re}(\Phi_h^0) , \\ H &= \cos\alpha \cdot \mathrm{Re}(\Phi_h^0) + \sin\alpha \cdot \mathrm{Re}(\Phi_\ell^0) . \end{aligned} \tag{18}$$

Now, since the physical Higgs masses are sensitive to the energy scale at which the quartic couplings λ_h and λ_ℓ are evaluated and since we do not apply the Higgs threshold corrections when solving the RGE, we need to estimate the errors on $m_h(\mu \sim m_h)$ and on $m_H(\mu \sim m_H)$. We do so by computing the difference between the masses obtained with $\lambda_{h/l}(\mu = 100 \text{ GeV})$ and with $\lambda_{h/l}(\mu = 400 \text{ GeV})$. We then find that the typical error is of $\mathcal{O}(\pm 10\%)$ for m_h and of $\mathcal{O}(\pm 20\%)$ for m_H.

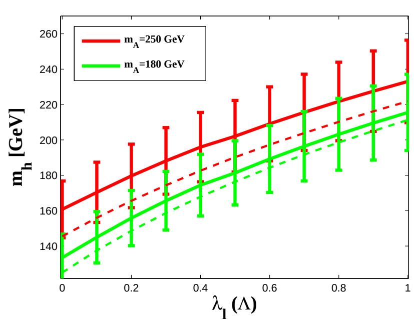

FIG. 1: *m_h as a function of $\lambda_\ell(\Lambda)$, for $\tan\beta = 0.7$, $m_{q'} = 400$ GeV and two representative values of m_A. The approximate analytic solutions are shown by solid lines and exact results (obtained from a full RGE analysis, see text) without errors by the dashed lines.*

[6] For an interesting recent DEWSB model with three generations and $\Lambda \gtrsim 10^{17}$ GeV, see [6].

As inputs we use m_A (recall that $\mu_{h\ell} = \sqrt{s_\beta c_\beta} m_A$), $\tan\beta$, $v = 246$ GeV, $m_{q'}$ (which sets the value of Λ, see Eq. 12) and $\lambda_\ell(\mu \sim \Lambda)$, while $\lambda_h(\mu)$ and $g_{q'}(\mu)$ are calculated from Eq. 11. In Fig. 1 we show the typical dependence of m_h on $\lambda_\ell(\Lambda)$, for $m_{q'} = 400$ GeV and for some representative values of $\tan\beta$ and m_A. We see that m_h decreases with $\lambda_\ell(\Lambda)$, so that the minimal m_h is obtained for $\lambda_\ell(\Lambda) = 0$, i.e., at the boundary below which the vacuum becomes unstable. Note that, for $\lambda_\ell(\Lambda) \to 0$, we find $\lambda_\ell(m_W) \sim \mathcal{O}(0.1)$ and, in particular, $\lambda_h(m_W) \gg \lambda_\ell(m_W)$. Thus, with $\lambda_\ell(\Lambda) \to 0$ and $\tan\beta \sim \mathcal{O}(1)$, we obtain (see Eqs. 13-16) $m_h \sim m_A/\sqrt{2}$ and $m_H \sim v\sqrt{\lambda_h/2}$.

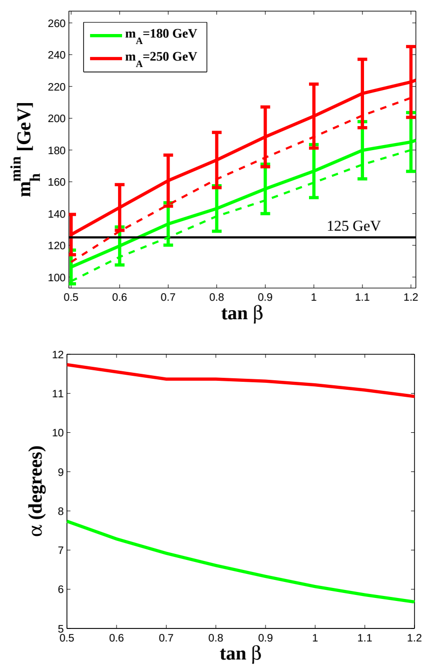

FIG. 2: *The minimal value of m_h (upper plot) and the corresponding values of the Higgs mixing angle α (lower plot), obtained by choosing $\lambda_\ell(\Lambda) = 0$ (see text), as a function of $\tan\beta$ for $m_A = 180$ and 250 GeV and $m_{q'} = 400$ GeV. See also caption to Fig. 1.*

In Fig. 2 we plot the minimal values/solutions for the lightest Higgs mass m_h (i.e., obtained with $\lambda_\ell(\Lambda) = 0$), as a function of $\tan\beta$, for $m_{q'} = 400$ GeV and $m_A = 180$, 200 and 250 GeV. We note that the dependence of our results on $m_{q'}$, or equivalently on $\Lambda(m_{q'})$, is negligible for values in the range 400 GeV$\lesssim m_{q'} \lesssim$600 GeV, as long as we choose the same boundary condition for $\lambda_\ell(\Lambda)$, i.e., at the compositeness scale.

We see that $m_h \sim 125$ GeV is obtained in the h4G2HDM with $\tan\beta \lesssim 0.7$ and $m_A \lesssim 250$ GeV. As shown above, this requires small values of $\lambda_\ell(\Lambda)$ and a small Higgs mixing angle of $\mathcal{O}(10^0)$ (also shown in Fig. 2), between the fundamental and the composite Higgs states. In particular, the light 125 GeV Higgs state in our model

is mostly the "fundamental" field, while the heavy CP-even Higgs is mostly a composite state, i.e., $h \sim \mathrm{Re}(\Phi_\ell^0)$ and $H \sim \mathrm{Re}(\Phi_h^0)$ (see Eq. 18). For the heavier Higgs we find that $m_H \sim 500 \pm 100$ GeV within the phenomenologically viable range of values of m_A and $\tan\beta$, which give $m_h \sim 125$ GeV.

Finally, we also depict in Figs. 1 and 2 the "exact" results, which are obtained from a full RGE analysis including the Yukawa couplings of the 4th generation leptons $\ell' = (\nu', \tau')$, of the top and of the bottom quarks, as well as the gauge couplings. In particular, assuming that ν' and τ' also couple to the auxiliary (composite) field Φ_h and using $m_{\ell'} = 200$ GeV at the EW-scale. Indeed, the observed slight shift from the approximate solutions is caused mainly by "turning on" the Yukawa couplings of the 4th generation leptons, and is within the estimated errors. We note, however, that the results are insensitive to the exact value of $m_{\ell'}$, so long as it is within the range $\mathcal{O}(m_W) < m_{\ell'} < \mathcal{O}(m_{q'})$.

To summarize, we have constructed a hybrid scenario for DEWSB, where both a fundamental-like Higgs field and a condensate of a strongly coupled heavy quark sector participate in EWSB. This yields a partly dynamical EWSB setup, and the resulting low-energy theory is a hybrid 2HDM with a "4th family" of heavy fermions, denoted here by h4G2HDM, in which one Higgs (mostly fundamental) couples only to the light SM fermions while the 2nd Higgs (mostly composite) couples only to the heavy (dynamical) 4th generation fermions. The proposed DEWSB framework results in a compositeness scale $\Lambda \sim \mathcal{O}(1)$ TeV and a phenomenologically viable low energy h4G2HDM setup, which closely resembles the recently proposed 4G2HDM of [15]. Thus, it is consistent with all current data [15, 16], including the recently measured 125 GeV Higgs signals [18], and may lead to interesting new signatures of the heavy quarks, such as the flavor changing decay $t' \to th$, which may be searched for in the present LHC data, see [17]. In particular, in spite of the heaviness of the dynamical quarks, $m_{q'} \sim \mathcal{O}(500)$ GeV, and the resulting low TeV-scale threshold for the strong dynamics, a viable Higgs candidate is obtained, which has a mass $m_h \sim 125$ GeV and properties *very similar* to the SM Higgs, but not identical, e.g., in the $\tau\tau$ and $\gamma\gamma$ channels, see [18]. This requires $\tan\beta \lesssim 1$, $m_A \sim m_{H^+} \lesssim 250$ GeV and $m_H \sim \mathcal{O}(500)$ GeV.

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